

Chiral condensate in nuclear matter with vacuum corrections

Wei-Zhou Jiang and Bao-An Li

Department of Physics, Texas A&M University-Commerce, Commerce, TX 75429, USA

Abstract

Within the relativistic Hartree approach using a Lagrangian with density-dependent parameters according to the Brown-Rho scaling law, it is found that the vacuum corrections from the nucleon Dirac sea soften the equation of state and favor the chiral symmetry restoration at high densities.

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1 Introduction

As hadrons can be viewed as excitations of the QCD vacuum, the vacuum structure that features varieties of quark and gluon condensates plays a crucial role in the hadron dynamics [1, 2, 3, 4, 5]. The chiral or quark condensate $\langle \bar{q}q \rangle$ serves as the order parameter for the spontaneous chiral symmetry breaking and reflects the nature of the QCD vacuum. For a chirally invariant vacuum, the chiral condensate vanishes. The non-vanishing chiral condensate plays a crucial role in generating the dynamical chiral symmetry breaking with the mass acquisition of constituent quarks. In free space, $\langle \bar{q}q \rangle_0 = -(225 \pm 25 \text{ MeV})^3$ according to the non-perturbative QCD approaches, for instance, see Refs. [4, 6, 7], whereas the perturbative vacuum is chirally symmetric for massless quarks. In the hot and/or dense medium, the partial restoration of chiral symmetry with reduced in-medium $\langle \bar{q}q \rangle$ is expected. The hadronic and electromagnetic signals for the partial restoration of chiral symmetry have been suggested in nucleus-nucleus reactions, for instance, see [3, 8, 9].

Due to the difficulties of the non-perturbative QCD, the in-medium chiral condensate may be derived alternatively using the QCD sum rules [10], and/or effective QCD models [1, 3, 11, 12, 13, 14]. The Nambu-Jona-Lasinio models and the linear σ models are among the well-known examples of the effective QCD models. Recently, the in-medium chiral condensate was also exploited in some new attempts [15, 16, 17]. A reduced in-medium chiral condensate was obtained in these models and approaches. However, an appropriate description of the nucleonic medium in terms of the quark degree of freedom is still very difficult. Interestingly,

a model-independent approach to evaluate the in-medium chiral condensate was developed by applying the Hellmann-Feynman (HF) theorem [11]. On the hadron level, the in-medium chiral condensate can be obtained from the derivative of the nuclear matter energy density with respect to the current quark mass. In fact, the in-medium chiral condensate has been evaluated using the nuclear matter energy density obtained from various many-body theories. For instance, the Dirac-Brueckner (DB) approach [18, 19], the relativistic mean field (RMF) models [20, 21, 22, 23], the quark meson coupling models [24, 25], and the effective field theory [26, 27], have all been used in studying the in-medium chiral condensate. Surprisingly, most of these studies have shown various hindrances towards the chiral symmetry restoration at high densities when nucleon-nucleon interactions are taken into account even in the modified quark-meson coupling model [28].

In principle, the non-vanishing chiral condensate arises from the vacuum of the non-perturbative QCD. In particular, the nucleonic vacuum can be considered as one of hierarchic constructions of the non-perturbative QCD vacuum. At finite density, the vacuum is modified by the medium. In dealing with the chiral condensate in nuclear medium, a proper treatment of the nucleonic vacuum is thus very important. However, so far very little attention has been paid to the vacuum corrections in evaluating the in-medium chiral condensate because usually only effective models are used. In this work, we explore the in-medium chiral condensate with the vacuum corrections using the Lagrangian constructed by us recently [29, 30] adopting the Brown-Rho (BR) scaling law for the partial restoration of chiral symmetry in nuclear matter [31, 32].

The paper is organized as follows. In Section 2, we renormalize our model with the Brown-Rho scaling in the Hartree approximation to attain the finite contribution from the Dirac sea. In section 3, we present the formalism for the in-medium chiral condensate using the HF theorem. Results on the Equation of State (EOS) and in-medium chiral condensate are discussed in Section 4. The summary is given in Section 5.

2 Relativistic Hartree model with BR scaling

In the present work, the model Lagrangian with density-dependent couplings and meson masses is written as [30]

$$\begin{aligned} \mathcal{L} = & \bar{\psi}[i\gamma_\mu\partial^\mu - M + g_\sigma^*\sigma - g_\omega^*\gamma_\mu\omega^\mu - g_\rho^*\gamma_\mu\tau_3b_0^\mu]\psi + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^{*2}\sigma^2) \\ & - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_\omega^{*2}\omega_\mu\omega^\mu - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + \frac{1}{2}m_\rho^{*2}b_{0\mu}b_0^\mu - \frac{1}{4}A_{\mu\nu}A^{\mu\nu} + \mathcal{L}_{CT}, \end{aligned} \quad (1)$$

where \mathcal{L}_{CT} is the counter term used for renormalization, ψ, σ, ω , and b_0 are the fields of the nucleon, scalar, vector, and isovector-vector mesons, with their masses $M, m_\sigma^*, m_\omega^*$, and m_ρ^* ,

respectively. The meson coupling constants and masses with asterisks denote their density dependence, given by the BR scaling [29, 30, 33]. The density dependence of parameters is described by the scaling functions that are the ratios of the in-medium parameters to those in the free space. We take the scaling function for meson masses as given in Ref. [29]

$$\Phi(\rho) = 1 - y\rho/\rho_0, \quad (2)$$

where the linear dependence on density is similar to the Nambu scaling [32] but with a much smaller coefficient $y \approx 0.1$ suggested by recent experiments [34, 35]. Here, the scaling for the nucleon is not considered since the nucleon mass already drops in the medium due to the coupling to the σ meson. The scaling functions for the coupling constants of scalar and vector mesons read

$$\Phi_\sigma(\rho) = \frac{1}{1 + x\rho/\rho_0}, \quad \Phi_\rho(\rho) = \frac{1 - y\rho/\rho_0}{1 + y_\rho\rho/\rho_0}, \quad \Phi_\omega(\rho) = \frac{1 - y\rho/\rho_0}{1 + y_\omega\rho/\rho_0}, \quad (3)$$

which are necessary for an appropriate description of nuclear matter properties [29].

The model renormalization is treated in the Hartree approximation. The Lagrangian for the counter terms reads

$$\mathcal{L}_{CT} = \alpha\sigma + \frac{1}{2!}\beta\sigma^2 + \frac{1}{3!}\gamma\sigma^3 + \frac{1}{4!}\lambda\sigma^4, \quad (4)$$

where the definition and meaning of the coefficients in \mathcal{L}_{CT} follow those given in Refs. [36, 37] but with the g_σ replaced by g_σ^* . These counter terms are used to renormalize the nucleon scalar self-energy, while no renormalization procedure is needed for the nucleon vector self-energy due to the nucleon current conservation. We use the procedure developed by Chin in Ref. [37] to renormalize the scalar self-energy and the energy density. In the following, we just stress the renormalization for the rearrangement term. The latter is essential for the thermodynamic consistency in deriving the pressure. In the Hartree approximation, the expectation value of the rearrangement term Σ_0^R is given by

$$\Sigma_0^R = \langle \frac{\partial \mathcal{L}}{\partial \rho} \rangle = -\rho^2 C_\omega \frac{\partial C_\omega}{\partial \rho} - \rho^2 \delta^2 C_\rho \frac{\partial C_\rho}{\partial \rho} - m_\sigma^* \sigma^2 \frac{\partial m_\sigma^*}{\partial \rho} + \Sigma^C, \quad (5)$$

and

$$\Sigma^C = \langle \bar{\psi}\psi \rangle \sigma \frac{\partial g_\sigma^*}{\partial \rho} + \frac{\partial \alpha}{\partial \rho} \sigma + \frac{1}{2!} \frac{\partial \beta}{\partial \rho} \sigma^2 + \frac{1}{3!} \frac{\partial \gamma}{\partial \rho} \sigma^3 + \frac{1}{4!} \frac{\partial \lambda}{\partial \rho} \sigma^4, \quad (6)$$

with

$$\langle \bar{\psi}\psi \rangle = -i \sum_{i=p,n} \int \frac{d^4 k}{(2\pi)^4} \text{tr} G(k) = \rho_S - i \sum_{i=p,n} \int \frac{d^4 k}{(2\pi)^4} \text{tr} G_F(k). \quad (7)$$

Here $C_\omega = g_\omega^*/m_\omega^*$, $C_\rho = g_\rho^*/m_\rho^*$, $\delta = (\rho_p - \rho_n)/\rho$ is the isospin asymmetry, ρ_S is the scalar density, and $G(k)$ is the nucleon propagator. After the renormalization, the Σ^C can be expressed as

$$\Sigma^C = \sigma \frac{\partial g_\sigma^*}{\partial \rho} \tilde{\rho}_S = \sigma \frac{\partial g_\sigma^*}{\partial \rho} \left[\rho_S - \sum_{i=p,n} \frac{2}{(2\pi)^2} \left(M^{*3} \ln \frac{M^*}{M} + \Sigma_S M^2 - \frac{5}{2} \Sigma_S^2 M + \frac{11}{6} \Sigma_S^3 \right) \right], \quad (8)$$

where the scalar nucleon self-energy $\Sigma_S = M - M^*$ with $M^* = M - g_\sigma^* \sigma = M - g_\sigma^* \tilde{\rho}_S / m_\sigma^{*2}$. The energy density and pressure read, respectively,

$$\mathcal{E} = \frac{1}{2}C_\omega^2 \rho^2 + \frac{1}{2}C_\rho^2 \rho^2 \delta^2 + \frac{1}{2}\tilde{C}_\sigma^2 (M^* - M)^2 + \sum_{i=p,n} \frac{2}{(2\pi)^3} \int_0^{k_{Fi}} d^3k E^* + \Delta\mathcal{E} \quad (9)$$

$$p = \frac{1}{2}C_\omega^2 \rho^2 + \frac{1}{2}C_\rho^2 \rho^2 \delta^2 - \frac{1}{2}\tilde{C}_\sigma^2 (M^* - M)^2 - \Sigma_0^R \rho + \frac{1}{3} \sum_{i=p,n} \frac{2}{(2\pi)^3} \int_0^{k_{Fi}} d^3k \frac{\mathbf{k}^2}{E^*} - \Delta\mathcal{E}, \quad (10)$$

where $\tilde{C}_\sigma = m_\sigma^* / g_\sigma^*$, and $\Delta\mathcal{E}$ is the finite vacuum correction [36, 37]

$$\Delta\mathcal{E} = -\frac{1}{4\pi^2} \left(M^{*4} \ln \frac{M^*}{M} + M^3 \Sigma_S - \frac{7}{2} M^2 \Sigma_S^2 + \frac{13}{3} M \Sigma_S^3 - \frac{25}{12} \Sigma_S^4 \right). \quad (11)$$

3 In-medium chiral condensate

The HF theorem transmits a parameter dependence of the Hamiltonian to that of the system energy. For the QCD Hamiltonian, one may divide it into a chirally invariant part and a part depending on the current quark mass, namely $\mathcal{H}_{QCD} = \mathcal{H}_0 + m_q \bar{q}q$. By taking m_q as the very parameter, the HF theorem reads

$$\langle \psi | \frac{\partial H_{QCD}}{\partial m_q} | \psi \rangle = \frac{\partial}{\partial m_q} \langle \psi | H_{QCD} | \psi \rangle. \quad (12)$$

Since the vacuum expectation value (VEV) $\langle 0 | H_{QCD} | 0 \rangle$ is subtracted to obtain the total energy of nuclear matter, the in-medium chiral condensate $\langle \bar{q}q \rangle_\rho$ is given as [11]

$$\langle \bar{q}q \rangle_\rho = \langle \bar{q}q \rangle_0 + \frac{\partial \mathcal{E}}{\partial m_q}. \quad (13)$$

Here, \mathcal{E} is given in Eq.(9), and can be written more generally as

$$\mathcal{E} = (M + E/A)\rho, \quad (14)$$

where E/A is the binding energy per nucleon. Both hadron masses and meson-nucleon coupling constants depend on the current quark mass. However, based on an analysis in the linear σ model, the derivative of the coupling constant with respect to the current quark mass is smaller than that of the hadron masses by a factor σ_N/M [19], with the pion-nucleon sigma term $\sigma_N \approx 45$ MeV. Thus, we neglect the dependence of the coupling constants on the current quark mass. The in-medium chiral condensate is then given by

$$\langle \bar{q}q \rangle_\rho = \langle \bar{q}q \rangle_0 + \left(\frac{dM}{dm_q} + \frac{\partial(E/A)}{\partial M} \frac{dM}{dm_q} \sum_{\sigma,\omega,\rho} \frac{\partial(E/A)}{\partial m_i^*} \frac{dm_i^*}{dm_q} \right) \rho. \quad (15)$$

The current quark mass derivative of the nucleon mass is expressed in terms of the σ_N as [11, 18]

$$\frac{dM}{dm_q} = \frac{\sigma_N}{m_q}. \quad (16)$$

For the vector (ω and ρ) mesons, one may resort to the constituent quark model and give the vector meson sigma term using the following relation

$$\frac{m_q}{m_V^*} \frac{dm_V^*}{dm_q} = \frac{m_q}{M} \frac{dM}{dm_q} = \frac{\sigma_N}{M}. \quad (17)$$

This approximation has been widely used in the literatures [11, 18, 19, 28]. The determination from the πN scattering gives $\sigma_N = 45 \pm 7$ MeV [38]. This central value can be reached in the constituent quark model when the quark confinement condition is imposed [13]. Recently, in Ref. [27] the two-pion exchange processes including the intermediate Δ isobar excitations instead of the surrogated σ -meson exchange were considered, and a hindered tendency toward the chiral symmetry restoration at high densities was still shown. In these chiral effective field theories [26, 27], the short-range Nucleon-Nucleon (NN) interactions only play a minor role on the change of the in-medium quark condensate, while the long- and intermediate-range NN interactions from one- and two-pion exchanges dominate the hindrance. This observation is instructive, though different from that in the one-boson-exchange approaches [18, 19, 20, 21, 22, 23, 24, 28]. However, the strict renormalization for the low-energy physics and contact interactions in these chiral effective field theories is nevertheless imperative to further judge whether the additional degrees of freedom are necessary for describing the short-range interactions. Though the progress of renormalization has been made recently [39, 40, 41], this is still quite challenging. In these chiral effective theories that were built upon the chiral perturbative theory, the vacuum is alienated in two folds: First, the pion as the Goldstone boson that resides in the coset space is not accompanied by other boson degrees of freedom that carry the chirally invariant vacuum in the conserved subgroup space. Second, in a non-relativistic treatment the nucleonic vacuum is out of touch. In fact, suitable renormalization scheme of the baryonic vacuum allows one to include consistently the vector mesons as explicit degrees of freedom in the Lorentz-invariant baryon chiral perturbative theory [42]. Since the chiral symmetry restoration is expected to occur at the short-range region with a chirally invariant vacuum, the present approach is of close relevance to dealing with the in-medium chiral condensate regarding the importance of the vacuum correction and the important role of short-range ω -meson exchange. As we focus on the vacuum correction, the relation (17), though phenomenological, can provide appropriate and referential information.

The situation with the scalar σ meson is more involved since it has a broad mass spectrum starting roughly above $2m_\pi$, as a result of the equivalence of a variety of two-pion exchanges.

As known in the linear σ model of Gell-Mann and Levy [43], the σ meson is regarded as the chiral partner of the Goldstone boson (π meson). In this model the chiral symmetry and its spontaneous breaking are well manifested. Although one can have some other choices to scale the sigma term for the σ meson [11, 18, 19, 28], we stress here that the symmetry of chirality is important and thus in our model that is designed to respect the chiral symmetry restoration at high densities we persist in using the sigma term for the σ meson obtained from the linear σ model, and at tree level it is [12, 19, 44]

$$\frac{m_q}{m_\sigma^*} \frac{dm_\sigma^*}{dm_q} = \frac{3}{2} \frac{\sigma_N}{M}. \quad (18)$$

The use of this relation is consistent with present relativistic Hartree approach which is also at tree level, though there can be high-order corrections to this relation. Here, the σ -meson mass is taken to be 590 MeV. It was obtained by fitting ground-state properties of many nuclei with the same scaling functions used in our previous work [30]. Using relations (2), (16), (17), (18) and the Gell-Mann-Oakes-Renner relation $m_\pi^2 f_\pi^2 = -m_q \langle \bar{q}q \rangle_0$ [6], the in-medium chiral condensate is eventually given by

$$\frac{\langle \bar{q}q \rangle_\rho}{\langle \bar{q}q \rangle_0} = 1 - \frac{\sigma_N}{m_\pi^2 f_\pi^2} \rho \left(1 + \frac{\partial(E/A)}{\partial M} + \frac{\partial(E/A)}{\partial m_\sigma^*} \frac{3m_\sigma}{2M} \Phi(\rho) - C_\omega^2 \frac{\rho}{M} - C_\rho^2 \frac{\rho \delta^2}{M} \right), \quad (19)$$

with the π -meson decay parameter $f_\pi = 93$ MeV and the π -meson mass $m_\pi = 138$ MeV. The first term in the bracket is the leading-order model-independent contribution. In this work, we focus on the chiral condensate in symmetric nuclear matter where the ρ meson does not contribute. We also do not consider the δ , π and η mesons, as they do not contribute to the symmetric nuclear matter either in the Hartree approximation. In fact, within the Dirac-Brueckner approach, their contributions to the in-medium condensate are about two orders of magnitude smaller than those of σ and ω mesons [18, 19]. Note that the role of the π meson in these relativistic one-boson-exchange theories is different from that in the effective field theories [26, 27], where the perturbative expansion with the unique π meson degrees of freedom results in the multi-pion exchanges [39].

4 Results and discussions

Before presenting our results on the in-medium chiral condensate, we discuss more about the determinations of model parameters. In our previous work [29, 30], besides the saturation properties of nuclear matter, namely, the incompressibility $\kappa = 230$ MeV at saturation density $\rho_0 = 0.16 \text{ fm}^{-3}$, we also use as a constraint the EOS of high density nuclear matter determined from analyzing the collective flow of relativistic heavy-ion collisions[45]. Here, we use the same

scaling functions of meson masses and meson-nucleon coupling constants as for the model SL1 constructed in Refs. [29, 30]. The model SL1, however, does not have the vacuum renormalization. The resulting new parameter set RSL1 and the corresponding saturation properties are tabulated in Table 1. Compared to the parameter set SL1, in the RSL1 the coupling constant of the ω meson is significantly reduced due to the vacuum contribution. Consistent with the finding in Ref. [37], the nucleon effective mass at saturation density increases significantly due to the vacuum correction. This means that the in-medium nucleon feels a scalar repulsion from the Dirac sea, and it prevents nucleons in the Fermi sea from falling to the fully-occupied Dirac sea by sustaining a larger mass gap.

Table 1: Parameter sets RSL1, RSL2 and SL1. The vacuum hadron masses are $M = 938\text{MeV}$, $m_\sigma = 590\text{MeV}$, $m_\omega = 783\text{MeV}$ and $m_\rho = 770\text{MeV}$. The coupling constants given here are those at zero density. The symmetry energy is fitted to 31.6MeV at saturation density.

	g_σ	g_ω	g_ρ	y	x	y_ω	y_ρ	$\kappa(\text{MeV})$	M^*/M
RSL1	8.7376	7.8476	4.0371	0.126	0.2831	-	-	230	0.788
RSL2	8.3034	6.9439	3.9878	0.126	0.2988	-0.0225	-	230	0.810
SL1	10.1937	10.4634	3.7875	0.126	0.234	-	-	230	0.679

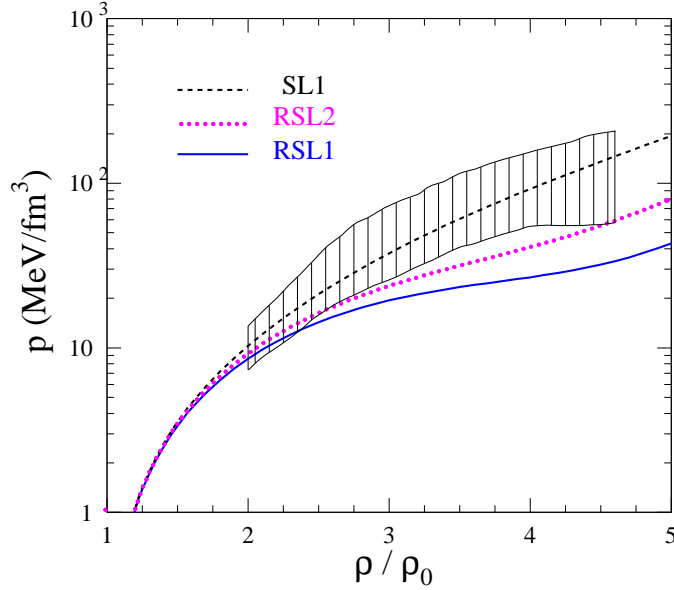


Figure 1: (Color online) The pressure for different models versus density in symmetric nuclear matter. The shaded region is given by experimental error bars[45].

Shown in Fig. 1 are the EOSs of symmetric nuclear matter for the RMF model SL1 and the relativistic Hartree models. The shaded area indicates the experimental constraint on the high-density EOS[45]. The model parametrization in the SL1 was selected such that the resulting EOS passes through the middle of the constraint. Comparing to the EOS with the SL1, it is seen that the pressure with the RSL1 is much reduced at high densities due to the vacuum correction. In order to bring the predicted pressure closer to the experimental region, we use another parameter set RSL2 that is adjusted to have a larger nucleon effective mass at saturation density, as listed in Table 1. As shown in Fig. 1, the moderate readjustment of parameters results in a significant increase of pressure at high densities where the rearrangement term plays an important role.

We have noticed that the energy density at high densities is dominated by the repulsion provided overwhelmingly by the ω meson. This is the main origin of the hindrance toward a chiral symmetry restoration observed using energy densities predicted by most meson-exchange many-body theories. In the simple RMF models, the repulsion increases linearly with density. After taking into account the higher-order correlations beyond the mean-field approximation within the DB approach, though the EOS at high densities was softened, the increasing tendency of the in-medium chiral condensate was not overturned [18, 19]. On the other hand, this may also be understandable since these approaches are not really expected to give a correct description of chiral condensate at least near the critical density for the chiral symmetry restoration. It was found very recently within the Brueckner approach with chiral limit that nucleon-nucleon interactions lead to a weaker reduction of the in-medium chiral condensate in the density region where the effective field theory is applicable [26], while little attention was paid to the saturation properties. The decreasing tendency of the in-medium chiral condensate can be obtained in RMF models by including the density-dependent coupling constants, namely the density-dependent ratios \tilde{C}_σ and C_ω [20]. In this work, instead of modifying the ratios \tilde{C}_σ and C_ω , we will investigate the role played by the vacuum corrections.

In Fig. 2, we show the in-medium chiral condensate evaluated using Eq.(19) with the RMF models NL3 [46] and SL1 and the relativistic Hartree models RSL1 and RSL2. Compared to the NL3 result, the in-medium chiral condensate with the SL1 is largely reduced at higher densities. This can be attributed to a much softened EOS with the SL1 [29], compared to that with the NL3 [45]. At high densities, however, the chiral condensate with the SL1 increases with density given the fact that the constant ratio C_ω in SL1 does not suppress the vector part of the energy density. The latter is quadratic in density, also see Eqs.(9) and (19). It is seen that the mass dropping given by the BR scaling is not sufficient to reduce the chiral condensate to zero due to the nucleon-nucleon interactions in nuclear matter.

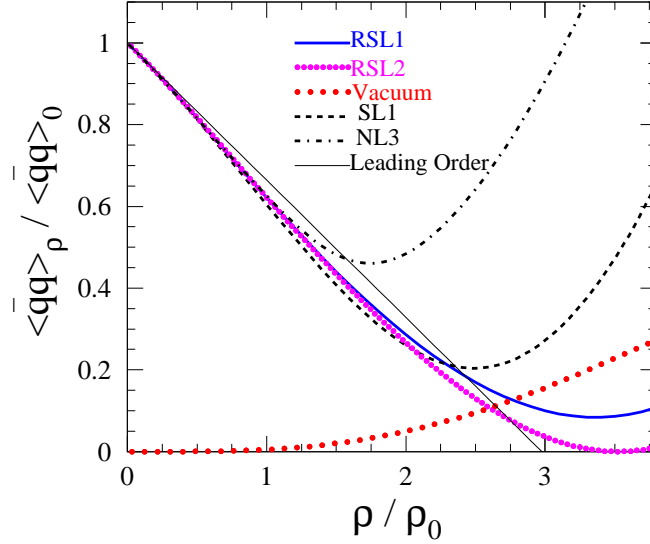


Figure 2: (Color online) In-medium chiral condensate for various models. The "vacuum" is calculated with the RSL2.

As the chiral condensate arises from the vacuum, the appropriate description of the Dirac vacuum contribution is imperative. The curve in the figure denoted by "vacuum" is calculated from the vacuum term $\Delta\mathcal{E}$ in Eq.(9). It is seen that the vacuum contribution to the in-medium chiral condensate becomes increasingly more important at higher densities. Moreover, the vacuum correction in the energy density also includes the part that modifies the nucleon effective mass. In the renormalized model RSL1 the chiral condensate further goes down with the increasing density. We note that the vacuum energy density $\Delta\mathcal{E}$ is positive and the scalar field σ that provides a big attraction is reduced due to the vacuum correction. Thus, the vacuum correction actually gives rise to a scalar repulsion that impairs the role of the vector repulsion provided by the ω meson. As a result, the much reduced g_ω with a consistent but smaller reduction of g_σ , subject to the given saturation properties, results in a softened EOS and brings down the chiral condensate at high densities in the RSL1, compared to that with the SL1.

Interestingly, when we use the parameter set RSL2 that is adjusted to have an EOS more close to the experimental constraint at high densities, see Fig. 1, the in-medium chiral condensate decreases further to zero by coincidence. As shown in Eqs.(9) and (10), the vacuum correction $\Delta\mathcal{E}$ contributes oppositely to the energy density and the pressure. This provides the intrinsic mechanism for the different changing tendencies in the energy density and the pressure, intrigu-

ing the shift going from the RSL1 to the RSL2 as observed in Figs. 1 and 2. Numerically, we find that a larger cancellation of contributions from the scalar and vector mesons in the RSL2, compared to the RSL1, is responsible for the big half of the reduction of the chiral condensate from the RSL1 to the RSL2, whereas the rest reduction arises directly from the variation of the vacuum correction. This further shows a direct link between the in-medium chiral condensate and the vacuum correction from the Dirac sea. With the vacuum correction, the nucleon effective mass drops more slowly in nuclear matter. As it was also stressed in the modified RMF models [20] and the Brueckner approach with the chiral limit [26], there is a decoupling between the reductions of the nucleon effective mass and the chiral condensate at high densities. Actually, this decoupling exists in most hadronic approaches where the in-medium chiral condensate increases whereas the nucleon effective mass drops at high densities.

In this work, we have concentrated on investigating effects of the vacuum corrections on the in-medium chiral condensate. However, if the chiral condensate does not drop sufficiently at high densities, other corrections, such as, high-order correlation contributions, may help further bring it down. Another simple way is to reduce the in-medium vector coupling constant. The expectation of vanishing chiral condensate at high densities would actually add some constraints on the in-medium scaling of the vector coupling constant in the framework of the BR scaling. It is also interesting to point out that the parameter used in the scaling function for the meson masses is consistent with that extracted from recent low-density experimental data, see Ref. [29]. However, it may be risky to apply the currently used scaling function having a linear density dependence of meson masses to high densities. In Ref. [32], different scalings were suggested for the regions below and above the saturation density. Moreover, one may consider the use of the nucleon mass scaling. Theoretical uncertainty may also exist in the approximations to determine the sigma terms for the vector and scalar mesons. Certainly, all these details deserve further investigations. Nevertheless, regardless of the very details and the uncertainties that ebb and flow, a definite fact is that the vacuum correction brings out a new mechanism favoring the reduction of the in-medium chiral condensate.

5 Summary

In summary, we have studied the vacuum corrections in the relativistic Hartree model using the density-dependent parameterizations according to the BR scaling. The in-medium chiral condensate have been investigated using the Hellmann-Feynman theorem together with phenomenological relations for the current quark mass derivative of hadron masses. Constrained by the saturation properties, the vector repulsion is reduced by the scalar repulsion arising from

the Dirac sea. This results in the softened EOS and the reduction of the in-medium chiral condensate at high densities. Also, the vacuum correction from the Dirac sea plays a direct and important role in generating the in-medium chiral condensate. The interplay between the vacuum corrections to the energy density and the pressure provides a mechanism favoring the chiral symmetry restoration in dense matter.

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References

- [1] V. Koch, Int. J. Mod. Phys. **E6**, 203 (1997); nucl-th/9706075.
- [2] C. M. Ko, V. Koch, G. Li, Annu. Rev. Nucl. Part. Sci. **47**, 505 (1997).
- [3] W. Cassing, and E.L. Bratkovskaya, Phys. Rep. **308**, 65 (1999).
- [4] E. O. Drukarev, M. O. Ryskin and V. A. Sadovnikova, Prog. Part. Nucl. Phys. **47**, 73 (2001).
- [5] G. E. Brown and M. Rho, Phys. Rep. **363**, 85 (2002).
- [6] M. Gell-Mann, R. Oakes and B. Renner, Phys. Rev. **175**, 2195 (1968).
- [7] J. Gasser and H. Leutwyler, Phys. Rep. **87**, 77 (1982).
- [8] C. M. Ko, G. Q. Li, J. Phys. **G22**, 1673 (1996).
- [9] B. Friman, W. Norenberg, V.D. Toneev, Eur. Phys. J. **A3**, 165 (1998).
- [10] For review, see, T. D. Cohen, R. J. Furnstahl, D. K. Grjegel and X. Jin, Prog. Part. Nucl. Phys. **35**, 221 (1995).
- [11] T.D. Cohen, R.J. Furnstahl and D.K. Griegel, Phys. Rev. **C45**, 1881 (1992).
- [12] G. Chanfray, M. Ericson, and P. A. M. Guichon, Phys. Rev. **C63**, 055202 (2001).

- [13] G. X. Peng, U. Lombardo, M. Loewe, and H. C. Chiang, Phys. Lett. B**548**, 189 (2002).
- [14] W. Z. Jiang, X.J. Qiu, Z.Y. Zhu and Z.J. He, Phys. Rev. C**65** (2002) 015210.
- [15] T. Mitsumori, N. Noda, H. Kouno, A. Hasegawa, M. Nakano, Phys. Rev. C**55**, 1577 (1997).
- [16] L. Chang, H. Chen, B. Wang, F. Y. Hou, Y. X. Liu, Phys. Lett. B**644**, 315 (2007).
- [17] Y. Kim, C. H. Lee, H.-U. Yee, Phys. Rev. D**77**, 085030 (2008).
- [18] G. Q. Li and C. M. Ko, Phys. Lett. B **338**, 118 (1994).
- [19] R. Brockmann and W. Weise, Phys. Lett. B **367**, 40 (1996).
- [20] A. Delfino, J. Dey, M. Dey and M. Malheiro, Phys. Lett. B **363**, 17 (1995).
- [21] A. Delfino, J. Dey, M. Dey, M. Malheiro, Z. Phys. C**71**, 507(1996).
- [22] M. Malheiro, M. Dey, A. Delfino, J. Dey, Phys. Rev. C**55**, 521 (1997).
- [23] M. Bagchi, M. Sinha, M. Dey, J. Dey, Phys. Lett. B**618**, 115 (2005).
- [24] K. Saito and A. W. Thomas, Phys. Rev. C**51**, 2757 (1995).
- [25] P. Wang, V. E. Lyubovitskij, T. Gutsche, A. Faessler, Phys. Rev. C**67**, 015210 (2003).
- [26] O. Plohl and C. Fuchs, Nucl. Phys. A **798**, 75 (2008).
- [27] N. Kaiser, P. de Homont and W. Weise, Phys. Rev. C **77**, 025204 (2008).
- [28] X. M. Jin and M. Malheiro, Mod. Phys. Lett. A **14**, 289 (1999).
- [29] W. Z. Jiang, B. A. Li and L. W. Chen, Phys. Lett. B**653**, 184 (2007).
- [30] W. Z. Jiang, B. A. Li and L. W. Chen, Phys. Rev. C**76**, 054314 (2007).
- [31] G. E. Brown, M. Rho, Phys. Rev. Lett. **66**, 2720 (1991).
- [32] G. E. Brown and M. Rho, Phys. Rep. **398**, 301 (2004).
- [33] C. Song, Phys. Rep. **347**, 289 (2001).
- [34] M. Naruki et al., Phys. Rev. Lett. **96**, 092301 (2006).
- [35] D. Trnka et al., Phys. Rev. Lett. **94**, 192303 (2005).
- [36] S. A. Chin, Phys. Lett. B**62**, 263 (1976).

- [37] S. A. Chin, Ann. Phys. **108**, 301 (1977).
- [38] J. Gasser, H. Leutwyler, and M. E. Sainio, Phys. Lett. B**253**, 252 (1991).
- [39] A. Nogga, R.G.E. Timmermans, U. van Kolck, Phys. Rev. C**72**, 054006 (2005).
- [40] M. P. Valderrama and E. R. Arriola, Phys. Rev. C**74**, 054001 (2006).
- [41] D.R. Entem, E. R. Arriola, M. P. Valderrama, R. Machleidt, Phys. Rev. C**77**, 044006 (2008).
- [42] T. Fuchs, M. R. Schindler, J. Gegelia, Stefan Scherer, Phys. Lett. B**575**, 11 (2003).
- [43] M. Gell-Mann and M. Levy, Nuovo Cim. **16**, 705 (1960).
- [44] M.C. Birse and J.A. McGovern, Phys. Lett. B**309**, 231 (1993).
- [45] P. Danielewicz, R. Lacey, W. G. Lynch, Science **298**, 1592, (2002).
- [46] G.A. Lalazissis, J.König, P.Ring, Phys. Rev. C**55**, 540 (1997).